

Preheating, Supersymmetry Breaking and Baryogenesis

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Fluctuations of scalar fields produced at the stage of preheating after inflation are so large that they can break supersymmetry much stronger than inflation itself. These fluctuations may lead to symmetry restoration along flat directions of the effective potential even in the theories where the usual high temperature corrections are exponentially suppressed. Our results show that nonthermal phase transitions after preheating may play a crucial role in the generation of the primordial baryon asymmetry by the Affleck-Dine mechanism. In particular, the baryon asymmetry may be generated at the very early stage of the evolution of the Universe, at the preheating era, and not when the Hubble parameter becomes of order the gravitino mass.

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In the low-energy minimal supersymmetric standard model (MSSM) there exist a large number of D -flat directions along which squark, slepton and Higgs fields may get expectation values. We collectively denote them by φ . In flat space at zero temperature exact supersymmetry guarantees that the effective potential $V(\varphi)$ along these D -flat directions vanishes to all orders in perturbation theory (besides the possible presence of nonrenormalizable terms in the superpotential). In the commonly studied supergravity scenario, supersymmetry breaking may take place in isolated hidden sectors [1] and then gets transferred to the other sectors by gravity. The typical curvature of D -flat directions resulting from this mechanism is $M_\varphi^2 \sim \frac{|F|^2}{M_P^2}$, where F is the vacuum expectation value (VEV) of the F -term breaking supersymmetry in the hidden sector. In order to generate soft masses of order of M_W in the matter sector, F is to be of order of $(M_W M_P)$ and sfermion and Higgs masses along flat directions turn out to be in the TeV range.

D -flat directions are extremely important in the Affleck-Dine (AD) scenario for baryogenesis [2] if large expectation values along flat vacua are present during the early stages of the evolving Universe. When calculating high temperature corrections to the effective potential at large φ one can often neglect the renormalizable interactions of the field φ with other fields, because the fields directly interacting with φ become too heavy to be excited by thermal effects. Therefore thermal effects do not push the field φ towards $\varphi = 0$ in the early universe if initially it was large enough.

For this reason, in the original version of the AD scenario [2,3] it was assumed that the initial value of the scalar field φ was different in different parts of the universe, as in the chaotic inflation scenario. However, recently it was pointed out that initial position of the field φ could be fixed by corrections to the effective potential which are proportional to the Hubble parameter H . Indeed, in generic supergravity theories moduli masses are

of order of the Hubble parameter H [4], [5]. One of the reasons is that the effective potential provides a nonzero energy density $V \sim |F|^2$ which breaks supersymmetry. If the vacuum energy dominates, the Hubble parameter is given by $H^2 = (V/3M_P^2)$ and therefore the curvature along the D -flat directions becomes $M_\varphi^2 = cH_I^2$, where c may be either positive or negative. (Here we use the reduce Planck mass $M_P \sim 2 \times 10^{18}$ GeV.) A similar contribution appears after inflation as well. For a mechanism of stronger supersymmetry violation see [6].

The most relevant part of the entire AD potential in presence of nonrenormalizable superpotentials of the type $\delta W = (\lambda/nM^{n-3})\varphi^n$ somewhat schematically can be represented as follows [4]:

$$V(\varphi) = c_1 m_{3/2}^2 |\varphi|^2 - c_2 H^2 |\varphi|^2 + \frac{A\lambda H\varphi^n e^{i n \theta} + h.c.}{n M^{n-3}} + \lambda^2 \frac{|\varphi|^{2n-2}}{M^{2n-6}}. \quad (1)$$

Here c_i are some positive constants which we suppose to be $O(1)$, $m_{3/2} = O(1)$ TeV is the gravitino mass, M is some large mass scale such as the GUT or Planck mass, n is some integer, which may take such values as 4, 6, etc., and the A -term violates the baryon or the lepton number and has a definite CP -violating phase relative to φ . This leads to symmetry breaking with $\varphi \sim (HM^{n-3}/\lambda)^{1/n-2}$. For definiteness we will take here $n = 4$ and $M = M_P$.

The main idea of ref. [4] is that in the early universe the field φ will be fixed at $\varphi \sim \sqrt{HM_P/\lambda}$. When H drops down to $m_{3/2}$, the terms proportional to H become smaller than the first term in (1), and the field φ begins oscillating near $\varphi = 0$. If the field φ can be associated with flat directions for squark-slepton fields, then the oscillations of the field φ near $\varphi = 0$ in the presence of the A -term of eq. (1) may produce baryon asymmetry of the universe.

The aim of this Letter is to show that there is a much stronger mechanism of supersymmetry breaking in the

early universe. This mechanism may lead to corrections to the quadratic terms in the effective potential of the AD field which are much greater than $H^2\varphi^2$ and to A -terms much greater than the one of eq. (1). As a result, under certain conditions the Affleck-Dine mechanism of baryon asymmetry generation may become quite different from the version proposed in [4].

The mechanism of supersymmetry breaking which we are going to discuss is related to particle production at the first stage of reheating after inflation. Indeed, Kofman, Linde and Starobinsky have recently pointed out that a broad parametric resonance, which under certain conditions occurs soon after the end of inflation, may lead to a very rapid decay of the inflaton field [7]. The inflaton energy is released in the form of inflaton decay products, whose occupation number is extremely large. They have energies much smaller than the temperature that would have been obtained by an instantaneous conversion of the inflaton energy density into radiation. Since it requires several scattering times for the low-energy decay products to form a thermal distribution, it is rather reasonable to consider the period in which most of the energy density of the Universe was in the form of the nonthermal quanta produced by inflaton decay as a separate cosmological era, dubbed as preheating to distinguish it from the subsequent stages of particle decay and thermalization which can be described by the techniques developed in [8]. Several aspects of the theory of explosive reheating have been studied in the case of slow-roll inflation [9] and first-order inflation [10] and it has been recently proposed that non-thermal production and decay of Grand Unified Theory bosons at the intermediate stage of preheating may generate the observed baryon asymmetry [11].

One of the most important consequences of the stage of preheating is the possibility of nonthermal phase transitions with symmetry restoration [12], [13]. These phase transitions appear due to extremely strong quantum corrections induced by particles produced at the stage of preheating. What is crucial for our considerations is that parametric resonance is a phenomenon peculiar of particles obeying Bose-Einstein statistics. Parametric resonant decay into fermions is very inefficient because of Pauli's exclusion principle. This means that during the preheating period the Universe is populated exclusively by a huge number of soft bosons, and the occupation numbers of bosons and fermions belonging to the supermultiplet coupled to the inflaton superfield are completely unbalanced. Supersymmetry is then strongly broken during the preheating era and large loop corrections may arise since the usual cancellation between diagrams involving bosons and fermions within the same supermultiplet is no longer operative. Therefore all results obtained in [12], [13] apply to the modification of the effective potential along the flat directions. As we will see, the curvature $V''(\varphi)$ during the preheating era often is much larger than the effective mass H^2 that D -flat directions acquire in the inflationary stage.

Let us consider, *e.g.*, chaotic inflation scenario [14] where the inflaton field ϕ couples to a complex χ -field,

$$V = \frac{M_\phi^2}{2} \phi^2 + g^2 \phi^2 |\chi^2|. \quad (2)$$

We take the inflaton mass $M_\phi \sim 10^{13}$ GeV in order for the density perturbations generated during the inflationary era to be consistent with COBE data. Inflation occurs during the slow rolling of the scalar field ϕ from its very large value. Then it oscillates with an initial amplitude $\phi_0 \sim M_P$. Within a dozen oscillations the initial energy $\rho_\phi \sim M_\phi^2 \phi_0^2$ is transferred through the interaction $g^2 \phi^2 \chi^2$ to *bosonic* χ -quanta in the regime of parametric resonance [7]. At the end of the broad parametric resonance the field ϕ drops down to $\phi_e \sim 10^{17}$ GeV. An exact number depends logarithmically on coupling constants; we will use $\phi_e \sim 10^{17}$ GeV $\sim 5 \times 10^{-2} M_P$ for our estimates. After this stage, the universe is expected to be filled up with χ -bosons with very large occupation numbers $n_k \sim g^{-2}$, with relatively small energy per particle, $E_\chi \sim \sqrt{g M_\phi \phi_e} \sim 0.2 \sqrt{g M_\phi M_P}$. It is especially important that the amplitude of field fluctuations produced at that stage is very large [7], [12],

$$\langle \chi^2 \rangle \sim 5 \times 10^{-2} g^{-1} M_\phi M_P. \quad (3)$$

If, for example, the energy of the inflaton field after preheating were instantly thermalized, we would obtain a much smaller value $\langle \chi^2 \rangle \sim 10^{-4} M_\phi M_P$. In realistic models thermalization typically takes a lot of time, and the value of $\langle \chi^2 \rangle$ after complete thermalization is many orders of magnitude smaller than (3).

Note, that large fluctuations (3) occur only in the bosonic sector of the theory, thus breaking supersymmetry at the quantum level. Anomalously large fluctuations of $\langle \chi^2 \rangle$ lead to specific nonthermal phase transitions in the early universe [12], [13]. It will be interesting to study possible consequences of this effect for supersymmetric theories with flat directions.

Let us make the natural assumption that the field φ labelling the D -flat direction couples to the generic supermultiplet χ by a renormalizable interaction of the form $\mathcal{L}_{\text{int}} = h^2 |\varphi|^2 |\chi|^2$ coming from some F -term in the potential. Reheating gives an additional contribution to the effective mass along the φ -direction :

$$\Delta M_\varphi^2 \sim 2 h^2 \langle \chi^2 \rangle \sim 10^{-1} \frac{h^2}{g} M_\phi M_P. \quad (4)$$

As a result, the curvature of the effective potential becomes large and positive, and symmetry rapidly restores for $10^{-1} \frac{h^2}{g} M_\phi M_P > c_2 H^2$. At the end of preheating in our model $H \sim 2 \times 10^{-2} M_\phi$. Thus, for $c_2 = O(1)$ symmetry becomes restored for $h^2 \gtrsim 2 \times 10^{-8} g$. In other words, supersymmetry breaking due to preheating is much greater than the supersymmetry breaking proportional to H unless the coupling constant h^2 is anomalously small. This leads to symmetry restoration in the AD model soon after preheating.

Let us check the applicability of our results. First of all, parametric resonance takes place for $g\phi > M_\phi$, where $\phi \sim 10^{-2}M_P$ at the end of preheating. This implies that $g > 10^{-4}$. Secondly, even though self-interactions of the χ -field do not terminate the resonance effect since particles remain inside the resonance shell, creation of quanta different from χ may remove the decay products of the inflaton away from the resonance shell. This leads us to consider two different possibilities. The resonance does not stop if scatterings are suppressed by kinematical reasons, *i.e.* if the non-thermal plasma mass of the final states is larger than the initial energy of the χ 's. If the final states are identified with the φ -quanta, this happens for $h > g$. Otherwise, scatterings occur, but are slow enough to not terminate the resonance if the interaction rate $\Gamma \sim n_\chi \sigma$, where σ denotes the scattering cross section, is smaller than the typical frequency of oscillations at the end of the preheating stage $\sim g\langle\chi^2\rangle^{1/2}$. Identifying again the final states with the φ 's, this condition translates into $h^2 \lesssim 10^2 g$, which is compatible with the condition for symmetry restoration discussed above.

There is another necessary condition in order to have broad resonance and a rapid production of χ -particles: the constant contribution $h|\varphi|$ to effective mass of the field χ induced by the condensate $|\varphi_0|$ should be smaller than the typical energy of the decay products E_χ . This translates into a bound on the coupling λ when taking into account the constraint on h discussed previously. The exact bound depends upon the choice of the index n and the mass M . For $n = 4$ and $M = M_P$ one has the condition $h^2 < 3g\lambda$. This condition simultaneously guarantees that the energy density of the AD field φ is smaller than the inflaton energy density at the end of the broad resonance. All these constraints are not very restrictive. For example, one may take $h = 2g$ to satisfy the condition $h > g$. Then one has the constraints $g > 10^{-8}$ and $\lambda > g$.

Let us now describe the dynamics of the φ -field during preheating. Under the conditions described above, the AD field φ rapidly rolls towards the origin $\varphi = 0$ and makes fast oscillations about it with an initial amplitude given by $\varphi \sim |\varphi_0|$. The frequency of the oscillations is of order of $(\Delta M_\varphi^2)^{1/2}$ and is much higher than the Hubble parameter at preheating $\sim 10^{-1} M_\phi$. The field is underdamped and relaxes to the origin after a few oscillations.

The crucial observation we would like to make here is that baryon asymmetry can be efficiently produced during this fast relaxation of the field to the origin. Our picture is similar to that of ref. [4], but there are considerable differences. According to [4], the baryon asymmetry is produced at very late times when the Hubble parameter H becomes of order of the gravitino mass $m_{3/2} \sim \text{TeV}$. At this epoch the field becomes underdamped and the CP -violating A -terms in the Lagrangian are comparable to the baryon-conserving ones. The field feels a torque from the A -terms and spirals from the initial point $\varphi = |\varphi_0| e^{i\theta}$ inward in the harmonic potential, producing

a baryon asymmetry $n_B/n_\varphi = \mathcal{O}(1)$.

In our case baryon number production occurs at very early stages, during the preheating era. This is possible because the same nonthermal effects generating a large curvature at the origin give rise also to sizable CP -violating terms along the AD flat direction. Indeed, under very general assumptions, one can expect the presence in the Lagrangian of nonrenormalizable CP -violating terms of the type $\alpha(\chi^2/M_P^{m-2})\varphi^m + \text{h.c.}$, where $m > 2$. During the preheating period large amplitudes of the χ -field give rise to the CP -violating term

$$\alpha \frac{\langle\chi^2\rangle}{mM_P^{m-2}} \varphi^m e^{im\theta_\chi} + \text{h.c.}, \quad (5)$$

For $m = 3$ we get a CP -violating term $\frac{A\varphi^3 e^{im\theta_\chi}}{3} + \text{h.c.}$ with $A \sim 10^{-1} g^{-1} \alpha M_\phi$.

Note that fluctuations $\langle\chi^2\rangle$ were absent during inflation; they increase very sharply (exponentially) during preheating, and therefore they begin strongly affecting the position of the scalar field only at the end of the broad parametric resonance, when they reach their maximal value. Thus, to the first approximation one has the following initial conditions for the field φ at the end of the broad resonance: $|\varphi_0| \sim \sqrt{HM_P/\lambda} \sim \sqrt{M_\phi M_P/10\lambda}$, $\dot{\varphi}_0 = 0$. As for the complex phase θ , its initial value is determined by the H -dependent A -term (1). In general, this phase is quite different from the phase determined by eq. (5). As a result, the field φ will spiral down to $\varphi = 0$ acquiring baryon charge density $n_B = 2|\varphi^2|\dot{\theta}$.

After the end of the stage of broad parametric resonance the value of $\langle\chi^2\rangle$ decreases approximately as $a^{-2} \sim t^{-1}$. As a result, both ΔM_φ and A after preheating decrease as t^{-1} , *i.e.* in the same way as the Hubble constant H in the radiation dominated universe after preheating. Therefore the equation describing the relaxation of the φ -field toward the origin

$$\ddot{\varphi} + 3H\dot{\varphi} + \Delta M_\varphi^2 \varphi + A(\varphi^*)^2 = 0 \quad (6)$$

takes the following form:

$$\ddot{\varphi} + \frac{3}{2t}\dot{\varphi} + \frac{1}{2gt}(h^2 M_P \varphi + \alpha \varphi^{*2}) = 0. \quad (7)$$

Before going further, let us compare the magnitudes of the last two terms for $|\varphi_0| \sim \sqrt{HM_P/\lambda} \sim \sqrt{M_\phi M_P/10\lambda}$. One can easily see that the B -violating term is greater than the mass term for $\alpha > 10^3 h^2 \sqrt{\lambda}$. If, *e.g.*, one takes $h \sim 10^{-7}$, $\lambda \sim 10^{-2}$, one finds that baryon number violation may be substantial even if the coefficient α in the A -term (5) is extremely small, $\alpha \sim 10^{-12}$.

It is convenient to introduce new variables $y = \varphi \sqrt{\lambda t/M_P}$ and $\tau = h \sqrt{2t M_P/g}$. In these variables eq. (7) simplifies:

$$y'' + y + \frac{\alpha \sqrt{2}}{\tau h \sqrt{g\lambda}} y^{*2} = 0. \quad (8)$$

In the new variables the motion of the field $y = |y|e^{i\theta}$ begins at $|y_0| = 1$, at the initial moment $\tau_0 \sim 10^4 h/\sqrt{g}$. The condition of vanishing initial velocity of the field φ translates into the condition $y'_0 = y_0/\tau_0$. To calculate the ratio n_B/s , where $n_B = 2|\varphi|^2\theta$, and s is the entropy density, we will introduce a fictitious reheating temperature $T_R \sim 10^{-2}\sqrt{M_\phi M_P} \sim 5 \times 10^{13}$ GeV. This is the temperature which would be reached by our system if thermalization would occur instantaneously after the end of the broad parametric resonance at $\phi_e \sim 10^{-1} M_P$. Even though thermalization may occur much later, this concept may be quite useful because at the radiation dominated stage the energy density of the universe decreases in such a way that at the moment when thermalization actually occurs the resulting temperature T will be equal to the redshifted value of the fictitious reheating temperature, $T \sim T_R \sqrt{\tau_0/\tau}$. This leads to the following expression for the baryon number $B = n_B/s$ soon after preheating:

$$B = \frac{n_B}{s} \sim \frac{n_B}{10^2 T_R^3} \sim 2 \times 10^{-2} |y|^2 \theta' \frac{h}{\lambda \sqrt{g}} . \quad (9)$$

We solved eq. (8) numerically for various values of parameters and calculated the ratio of n_B to the entropy density s . We have found that this ratio oscillates and approaches a constant at large times. For $\frac{\alpha}{h\sqrt{g\lambda}} < 1$ and $h > 10^{-4}\sqrt{g}$ the typical value of the baryon asymmetry is given by

$$B = \frac{n_B}{s} \sim 10^{-2} \frac{\alpha}{g\lambda\sqrt{\lambda}} f(\tau_0) . \quad (10)$$

Here $f(\tau_0)$ is a certain function of $\tau_0 \sim 10^4 h/\sqrt{g}$: $f(1) \sim 1$, $f(10) \sim 0.1$, $f(100) \sim 0.05$.

Validity of this result depends on details of the theory; typically it gives a reliable estimate of the baryon asymmetry only for $B \ll 1$ (which is what we need), but even in this case some care should be taken. For example, one can show that for $h^2 < g\lambda$ the temperature T remains much greater than the time-dependent effective mass of the field φ until this mass approaches the constant value $O(m_{3/2})$. In this regime quarks acquire large effective mass $O(T)$, and the field φ cannot decay and transfer its baryon asymmetry to fermions until temperature drops to $O(m_{3/2})$. This may lead to some corrections to eq. (10) [3]. Also, the simple scaling rules used in the derivation of eq. (10) may break if, e.g., at some intermediate stage the universe becomes dominated by nonrelativistic particles.

A detailed investigation of baryon asymmetry production in realistic theories including all of the effects mentioned above should become a subject of a separate investigation. The main purpose of our paper was to show that parametric resonance and nonthermal phase transitions found in [7,12] may lead to strong supersymmetry breaking in the early universe and to considerable modifications in the theory of baryogenesis in supersymmetric

models. We have found that if the inflaton field ϕ couples in a renormalizable way to χ -bosons, which are weakly coupled to the AD field φ , then the effect of parametric resonance may induce very large masses for particles corresponding to flat directions of the AD potential. The same effect may induce large terms violating baryon conservation (5). As a result, one can obtain large baryon asymmetry even in the models where all relevant coupling constants are extremely small.

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